

Slow light with three-level atoms in metamaterial waveguides

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Metamaterial is promising for enhancing the capability of plasmonic devices. We consider a cylindrical waveguide with three-level Λ atoms embedded in the dielectric core. By comparing metal cladding vs metamaterial cladding of a waveguide with Λ atoms in the core, we show that, for a fixed amount of slowing of light due to electromagnetically induced transparency, the metamaterial cladding outperforms in terms of the inherent loss.

I. INTRODUCTION

Plasmonic devices show promise in increasing the speed of electronic devices and networks while still meeting the size requirements of modern electronic devices due to their ability to confine light to sub-wavelength scales [1]. Combining such devices with optical phenomena, such as slow light, could yield an array of new optical devices to augment or replace existing technology. The feasibility of using nonlinear metamaterials for slow light has been considered [2, 3], and a scheme for photon-echo quantum memory for surface plasmon-polaritons on a metamaterial interface has been proposed [4]. Controllable slow light is useful for optical delay lines, optical buffers [5] and enhanced nonlinear interactions [6]; however, achieving highly-confined slow light with weak fields is impractical due to the high losses suffered as a result of the plasmonic confinement of light [7]. There exists a scheme [8, 9] using a flat interface metamaterial-dielectric waveguide for slowing highly-confined light through electromagnetically induced transparency (EIT) [10], while using metamaterial to minimize losses. EIT with weak fields is interesting in its own right for fundamental reasons, to distinguish it from Autler-Townes splitting, and also for applications to weak-field sensing [11]. However, the scheme has the limitation that field confinement in the waveguide is only in one transverse spatial dimension, which allows propagating fields to diverge.

We propose using a cylindrical waveguide structure as shown in Fig. 1, rather than a flat interface, to inherently provide confinement in both transverse directions and prevent field divergence. The waveguide is composed of a dielectric core and a metamaterial cladding that surrounds the core. Three-level Λ atoms (Fig. 2) are homogeneously embedded throughout the core of the waveguide and are driven by a pump field in a low-loss surface mode [9], thereby enabling EIT for slowing the signal field in another low-loss surface mode. To prepare EIT the atoms are initially strongly pumped to depopulate the $|s\rangle$ level and relax to the $|g\rangle$ state. This process creates Fano interference in the $|e\rangle \leftrightarrow |g\rangle$ transition resulting in a transparency window at the $|g\rangle \leftrightarrow |e\rangle$ transition frequency. As a result of the refractive index change, the group velocity of the signal pulse is reduced when traveling through the transparent medium. After preparing for the EIT, the intensity of the pump field is lowered because the weaker the pump field is the more reduction of group velocity of the signal field. Hence, coherently controlled slow light is achievable by tuning the intensity of the pump field. Furthermore, by comparing the metamaterial-cladded waveguide to a metal-cladded one with the same permittivity, we show that a low-loss surface mode [12] of the metamaterial-cladded guide enables

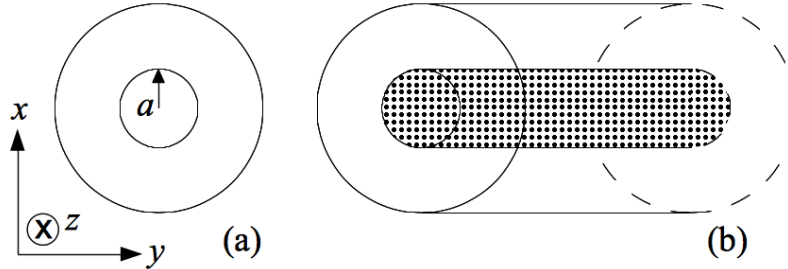


FIG. 1. (a) The cylindrical waveguide cross section. (b) A schematic of the metamaterial-dielectric waveguide embedded with three-level Λ atoms (black dots) in the core. The outer layer (cladding) is a metamaterial whereas the core is a dielectric.

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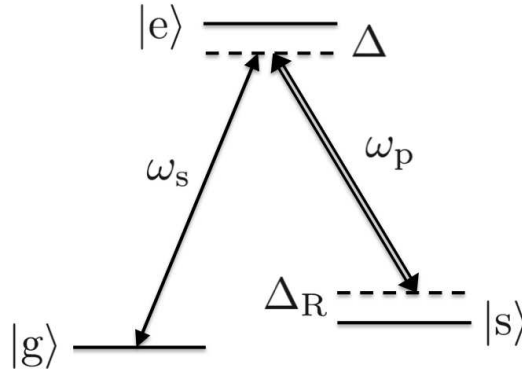


FIG. 2. A schematic of the three-level Λ atom. The ground state is $|g\rangle$, the excited state is $|e\rangle$ and $|s\rangle$ is a metastable state. The frequency of the pump field is ω_p and of the signal field is ω_s , and Δ and Δ_R are detunings.

the same degree of slowing of light but with reduced losses. Our analysis is important to fabricating coherently-controlled low-loss slow light devices with current metamaterial technology and we provide conditions necessary for the existence of the low-loss surface mode for slowing light in the waveguide.

In section II, we develop the theory for field propagation in the cylindrical metamaterial-cladded waveguides without the doping of the three-level atoms. Then we extend the theory to allow for doping the atoms in the core of the waveguide. In section III, we describe the numerical method we use to solve Maxwell's equations in the doped waveguide and present results obtained from the solutions. A discussion of our scheme is given in section IV and a summary of the work is given in section V.

II. THEORY

Metamaterials are designed with artificial structure, giving them electromagnetic responses different to those of their constituent materials. We consider a fishnet-type metamaterial as they have been well studied, both experimentally and theoretically [13–16]. Metamaterials of this design are inherently lossy so the permittivity and permeability expressions must necessarily include dissipation terms. The permittivity of the fishnet metamaterial is described by the Drude model [17],

$$\frac{\epsilon_{MM}(\omega)}{\epsilon_0} = \epsilon_b - \frac{\omega_e^2}{\omega(\omega + i\Gamma_e)}, \quad (1)$$

with ϵ_0 the permittivity of free space, ϵ_b the background permittivity, ω_e the plasma frequency of the metal, and Γ_e the electric damping term. The permeability, which is in practice achieved through the structure of the metamaterial, is described by the relation [18]

$$\frac{\mu_{MM}(\omega)}{\mu_0} = \mu_b + \frac{F\omega^2}{\omega_0^2 - \omega(\omega + i\Gamma_m)}, \quad (2)$$

with μ_0 the permeability of free space, μ_b the background permeability, F a geometric parameter, ω_0 the resonance frequency, and Γ_m the magnetic damping term. Equations (1) and (2) describe the electromagnetic properties of the metamaterial cladding and are used, along with a constant permittivity for the core, to determine the modes of the undoped metamaterial-dielectric waveguide.

After preparation for EIT, the $|s\rangle$ state is not populated and the atoms are no longer excited by the pump field, which means the pump field effectively travels in the modes of the undoped waveguide. On the other hand, the signal field can excite the atoms, so the signal field travels in the modes of the doped waveguide. In the following, we develop the theory for finding the modes in the undoped waveguide before finding the modes of the doped waveguide.

The guided modes in a cylindrical waveguide are transverse electric (TE), transverse magnetic (TM) or a combination of the two (denoted HE or EH) [19]. The TM and TE modes are those with field components $H_z = 0$ and $E_z = 0$ respectively, with z the propagation direction. We restrict our analysis to TM modes as they are invariant in azimuthal angle ϕ , which allows us to simplify the calculations. In general, the analysis can be done with any choice of modes.

For the TM modes of the metamaterial-dielectric waveguide, the electric fields have the form

$$\tilde{\mathbf{E}}(r, \phi, z, t) = \mathbf{E}(r, \phi) e^{i(\tilde{\beta}z - \omega t)} + \text{c.c.}, \quad (3)$$

with c.c. indicating the complex conjugate, and $\tilde{\beta} = \beta + i\alpha$ the complex propagation constants for the allowed modes of the waveguide, which are obtained by solving the dispersion relation. Using the definitions

$$J'_m(a\kappa) := \left. \frac{\partial J_m(r\kappa)}{\partial r} \right|_{r=a} \quad (4)$$

and

$$K'_m(a\gamma) := \left. \frac{\partial K_m(r\gamma)}{\partial r} \right|_{r=a}, \quad (5)$$

with J_m and K_m the Bessel function and modified Bessel function respectively, the dispersion relation for TM modes in an undoped cylindrical metamaterial-dielectric waveguide is [19]

$$\frac{\epsilon_d}{\kappa^2} \frac{J'_m(a\kappa)}{J_m(a\kappa)} + \frac{\epsilon_{\text{MM}}}{\gamma^2} \frac{K'_m(a\gamma)}{K_m(a\gamma)} = 0, \quad (6)$$

with

$$\gamma := \sqrt{\tilde{\beta}^2 - \omega^2 \epsilon_{\text{MM}} \mu_{\text{MM}}}, \quad \kappa := \sqrt{\omega^2 \epsilon_d \mu_d - \tilde{\beta}^2}, \quad (7)$$

ϵ_d the permittivity of the dielectric core, and a the core radius. With this dispersion relation, we can find the low loss TM surface mode for the pump field as detailed in [12].

The susceptibility of the undoped waveguide core is simply the dielectric susceptibility,

$$\chi_d = \epsilon_d / \epsilon_0 - 1. \quad (8)$$

When the core is doped with the three-level atoms, the susceptibility of the core becomes the sum $\chi_d + \chi(\omega, r)$, which alters the guided modes near the EIT resonance frequency from those of the undoped guide. To find the low loss TM surface mode of the signal field, we must solve the wave equation with the susceptibility of the doped waveguide core.

The susceptibility of the pumped three-level Λ atoms under EIT is [10]

$$\chi(\omega, r) = \frac{2a_0 c}{n_d \omega} \frac{i\gamma_{\text{eg}}}{\gamma_{\text{eg}} - i\Delta + |\Omega_p(r)|^2 (\gamma_{\text{sg}} - i\Delta_R)^{-1}}, \quad (9)$$

with c the speed of light in vacuum, $n_d = c\sqrt{\epsilon_d \mu_d}$ the refractive index of the dielectric core with μ_d the permeability of the dielectric core, γ_{ij} the decay rate from level i to level j and

$$a_0 := \frac{3\pi c^2}{n_d^2 \omega^2} \rho_a \approx \frac{3\pi c^2}{n_d^2 \omega_{\text{eg}}^2} \rho_a \quad (10)$$

with ρ_a the number density of three-level atoms. The approximation in Eq. (10) is valid near the EIT resonance frequency. Defining ω_{ij} as the transition frequency for the $|i\rangle \leftrightarrow |j\rangle$ transition the detunings are $\Delta = \omega_{\text{eg}} - \omega_s$, $\Delta_R := \Delta - \omega_{\text{es}} + \omega_p$. The pump is set on resonance with the $|s\rangle \leftrightarrow |e\rangle$ transition such that $\Delta_R = \Delta$. The Rabi frequency of the pump is

$$\Omega_p(r) = \frac{1}{\hbar} \mathbf{d} \cdot \mathbf{E}_p(r), \quad (11)$$

with \mathbf{d} the dipole moment of the $|s\rangle \leftrightarrow |e\rangle$ transition.

Equations (9) and (11) have r dependence but no ϕ or z dependence. This is because the pump is in a TM mode, which has a radial dependence but no azimuthal dependence, implying ϕ invariance. Furthermore, the pump is approximated as non-depleting, which implies z invariance. This approximation is valid for short propagation lengths, such that the pump does not deplete considerably. The results presented in this paper are based on this approximation.

For longer propagation lengths the pump does deplete, which reduces the Rabi frequency Ω_p and leads to a further reduction in the group velocity. In this case, one can treat the whole waveguide as a series of concatenated short

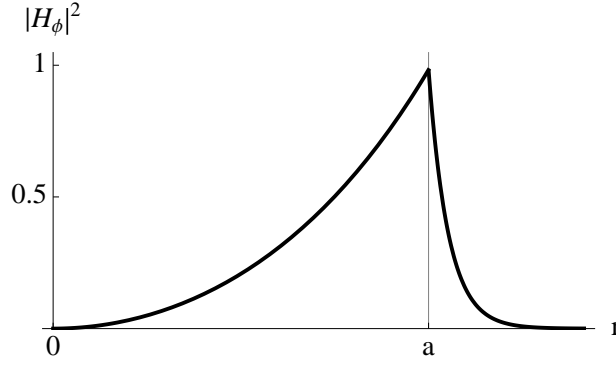


FIG. 3. A plot of $|H_\phi|^2$ of a surface mode showing the transverse profile in arbitrary units. The thin vertical line indicates the interface between the core and cladding regions.

sections along the propagation direction, such that the pump intensity in each section is approximately constant. This treatment allows one to determine the group velocity of the signal in each section and calculate the overall delay of the signal in the whole waveguide.

Inside the waveguide core, the pump field expression is given by

$$\mathbf{E}_p(r) = A \left(J_0(\kappa_p r) \hat{\mathbf{z}} - i \frac{\tilde{\beta}_p}{\kappa_p^2} J_1(\kappa_p r) \hat{\mathbf{r}} \right) \quad (12)$$

with $\hat{\mathbf{z}}$ and $\hat{\mathbf{r}}$ unit vectors in the propagation and radial directions respectively, A a constant that determines the amplitude of the pump, $\tilde{\beta}_p$ the complex propagation constant for the pump mode, and

$$\kappa_p := \sqrt{\omega_{\text{es}}^2 \epsilon_d \mu_d - \tilde{\beta}_p^2}. \quad (13)$$

We define $\eta := Ad$, which is a parameter adjusted by the amplitude of the pump, with $d = \mathbf{d} \cdot \hat{\mathbf{z}} = \mathbf{d} \cdot \hat{\mathbf{r}}$ such that,

$$\Omega_p(r) = \frac{\eta}{\hbar} \left(J_0(\kappa_p r) - i \frac{\tilde{\beta}_p}{\kappa_p^2} J_1(\kappa_p r) \right). \quad (14)$$

The permittivity of the doped core depends on the pump field. The pump field is in a TM surface mode (see Fig. 3), so the intensity varies steeply in the radial direction; thus, the signal field experiences a steeply graded refractive index in the core. As a result, within the core, the scalar wave equation used to solve for the modes of the undoped guide is no longer valid for the modes of the doped guide. In the cladding region, however, the permittivity and permeability are not spatially dependent, so we can still find the fields in this region using the scalar wave equation.

To solve for the propagation constants and the fields in the doped guide, we must begin with the vector wave equation

$$\nabla^2 \mathbf{E} + \nabla \left(\frac{1}{\epsilon_{\text{eff}}(\omega, r)} \mathbf{E} \cdot \nabla \epsilon_{\text{eff}}(\omega, r) \right) + k_0^2 \epsilon_{\text{eff}}(\omega, r) \mathbf{E} = 0, \quad (15)$$

with ∇^2 the Laplacian and

$$\epsilon_{\text{eff}}(\omega, r) = \begin{cases} \epsilon_0 (1 + \chi_d + \chi(\omega, r)), & r < a, \\ \epsilon_{\text{MM}}(\omega), & r \geq a, \end{cases} \quad (16)$$

the effective permittivity of the waveguide with three-level Λ -atoms embedded in the core. As we are restricting our analysis to TM modes for both the pump and signal fields, Eq. (15) can be decoupled into scalar equations with the z component taking the form

$$\frac{d^2}{dr^2} E_z + \left(\frac{1}{r} - \frac{\tilde{\beta}^2}{\kappa_{\text{eff}}^2(\omega, r) \epsilon_{\text{eff}}(\omega, r)} \frac{d}{dr} \epsilon_{\text{eff}}(\omega, r) \right) \frac{d}{dr} E_z + \kappa_{\text{eff}}^2(\omega, r) E_z = 0, \quad (17)$$

with

$$\kappa_{\text{eff}}(\omega, r) := \begin{cases} \sqrt{\omega^2 \mu_0 \epsilon_{\text{eff}}(\omega, r) - \tilde{\beta}^2}, & r < a, \\ i\gamma, & r \geq a. \end{cases} \quad (18)$$

Equation (17) is not generally analytically solvable, except at certain frequencies, so we numerically determine the dispersion and attenuation of the low-loss surface mode near the EIT resonance frequency.

III. METHOD AND RESULTS

There are two main problems with solving Eq. (17). First, $\tilde{\beta}$ is complex, meaning it has two free parameters so a standard shooting method cannot be employed [20]. The second problem is that the fitness function,

$$f = \left| \frac{H_{\phi, \text{core}} - H_{\phi, \text{clad}}}{H_{\phi, \text{clad}}} \right|, \quad (19)$$

with $H_{\phi, \text{core}}$ and $H_{\phi, \text{clad}}$ the ϕ component of the magnetic field in the core and cladding regions respectively, is not convex, so implementing a hill descent algorithm is not sufficient to determine if a particular value of $\tilde{\beta}$ is a solution. To circumvent these problems we draw inspiration from the shooting method and employ a trial method to solve Eq. (17) for $E_z(r, \omega)$.

We need to find $E_z(r, \omega)$ for a range of frequencies of interest. We begin at a frequency where an exact solution for a TM surface mode is found. To find the solution at an adjacent frequency, we slightly deviate the known $\tilde{\beta}$ at the former frequency to obtain a trial value for $\tilde{\beta}$ and numerically solve Eq. (17) for $E_z(r, \omega)$; other components of the field are calculated from this trial $E_z(r, \omega)$. To test whether the trial $\tilde{\beta}$ is acceptable, we check if the corresponding trial $E_z(r, \omega)$ satisfies the boundary condition at the core-cladding interface to within a tolerance. Mathematically, the acceptance criterion is $f < 10^{-9}$. If the trial $E_z(r, \omega)$ does not satisfy the boundary condition, a different trial $\tilde{\beta}$ is chosen by slightly deviating the known $\tilde{\beta}$ at the former frequency and the search is repeated. When an acceptable $\tilde{\beta}$ is found, we repeat the search for the next frequency until $E_z(r, \omega)$ of the TM surface mode is determined for a range of frequencies of interest.

We apply the same metamaterial parameters presented in our previous work [12] to our calculations in this work. Hence, the TM surface mode of the undoped guide found in [12] is the mode for the pump of our scheme. The metamaterial parameters are shown in Table I. The core parameters are $\epsilon_d = 1.3\epsilon_0$, $\mu_d = \mu_0$ and the core radius $a = 4\pi c/\omega_e$. For the three-level atoms embedded in the core, we set the number density $\rho_a = 1.26 \times 10^{21} \text{m}^{-3}$ so that the resulting permittivity remains positive. The decay rate of the excited state is $\gamma_{\text{eg}} = 10^5 \text{s}^{-1}$ and that of the hyperfine state is $\gamma_{\text{sg}} = 10^{-2} \text{s}^{-1}$; both of which are for Pr^{3+} [21]. The mean signal frequency is $\omega_s = \omega_{\text{eg}} = 0.409\omega_e$ and the mean pump frequency is $\omega_p = \omega_{\text{es}} = 0.41\omega_e$. The group velocity of the pump mode $v_p = 0.47c$.

TABLE I. Parameter values for the metamaterial cladding.

Parameter	Value
ϵ_b	1
μ_b	1
ω_e	$1.37 \times 10^{16} \text{s}^{-1}$
Γ_e	$2.73 \times 10^{13} \text{s}^{-1}$
Γ_m	Γ_e
ω_0	$0.2\omega_e$
F	0.5

Using the method and parameters outlined above, we have calculated the propagation constants of the signal and pump TM surface modes of our scheme. With the complex propagation constant $\tilde{\beta} = \beta + i\alpha$ for a range of frequencies, we have calculated the group velocity of the signal field using $v_g = (d\beta/d\omega)^{-1}$. Figure 4 shows the group velocity of the signal field surface mode in both the metal- and metamaterial-cladded waveguides for a range of η . The plot shows that the group velocity of the signal for a metamaterial-cladded guide is proportional to η^2 , and as $\eta \propto |\Omega(r)|$, this agrees with EIT theory [10].

Using the imaginary part of the propagation constant, we have obtained the attenuation curves for both the metamaterial- and metal-cladded waveguides as a function of detuning Δ as shown in Fig. 5. At the EIT resonance

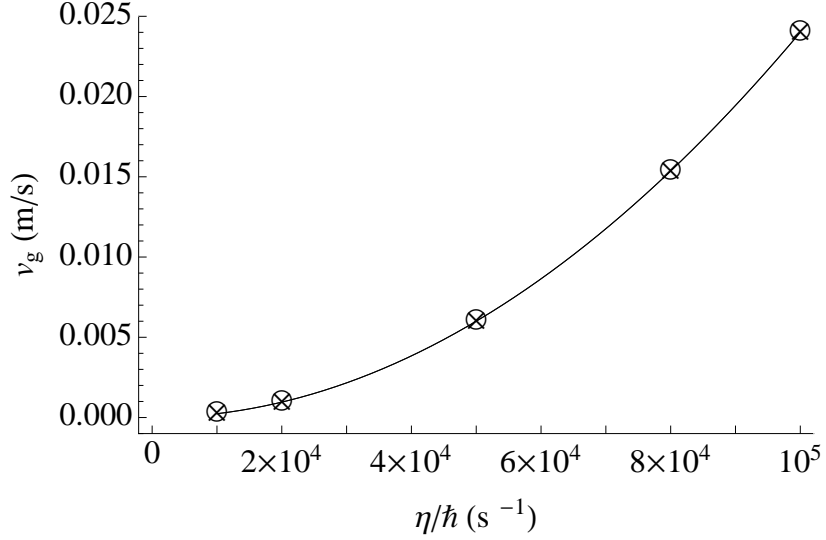


FIG. 4. A plot of the numerically calculated group velocity at $\Delta = 0$ as a function of pump intensity for a metamaterial-cladded guide (\circ) with $\Gamma_m = \Gamma_e/3$ and a metal-cladded guide (\times). The solid curve is a least squares quadratic fit.

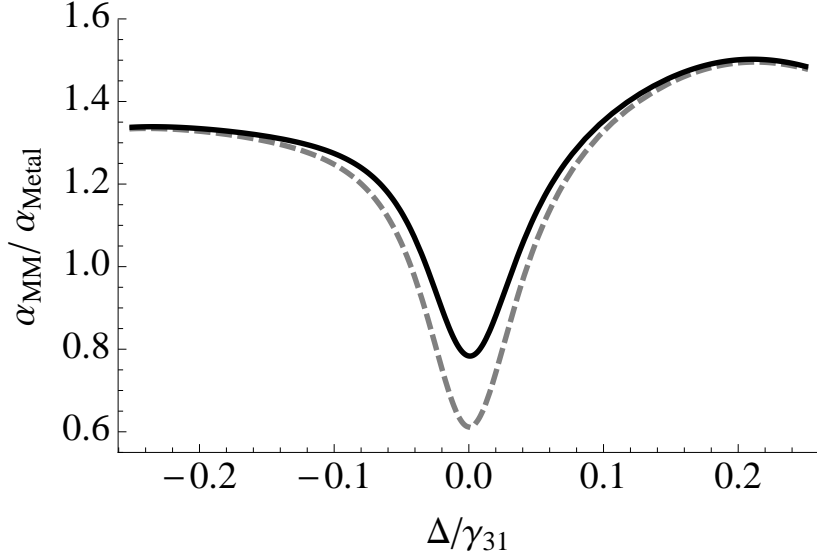


FIG. 5. Plots of the relative attenuation of the surface mode of the metamaterial guide to that of a metal guide ($\alpha_{MM}/\alpha_{Metal}$) with $\eta/\hbar = 5 \times 10^4 \text{ s}^{-1}$, for $\Gamma_m = \Gamma_e/3$ (solid) and $\Gamma_m = \Gamma_e/100$ (dashed).

(i. e. zero detuning), the metamaterial-dielectric guide provides nearly a 20% reduction in attenuation over a metal-dielectric guide for $\Gamma_m = \Gamma_e/3$ (solid line). For $\Gamma_m = \Gamma_e/100$ (dashed line), the attenuation reduction improves to near 40% over a metal-dielectric guide. When $\Gamma_m = \Gamma_e$ the attenuation of a metamaterial-dielectric guide is comparable to a metal-dielectric guide. For layered structures, such as the fishnet design, previous work [18] has shown that $\Gamma_m \leq \Gamma_e$ and a value of $\Gamma_m = \Gamma_e/3$ should be feasible with current metamaterial technology [22]. Hence, our cylindrical metamaterial-dielectric waveguide is capable of supporting a low loss TM surface mode.

To understand why the attenuation of the low-loss surface mode does not approach zero near ω_s , as it does for the flat interface [8], we plot the fraction of the total energy of the mode that resides in both the core and cladding. Figure 6 shows the fractional energy, ξ , of the TM surface mode of the cylindrical metamaterial-dielectric guide as a function of frequency. For a given mode the fractional energy is the ratio of the energy in one part of the guide, the core for instance, to the total energy of the mode. The solid line is for the fractional energy in the core whereas the dashed line is for the cladding. The fraction of energy in the dielectric does not reach one and some of the signal field penetrates into the metamaterial, leading to a non-zero attenuation.

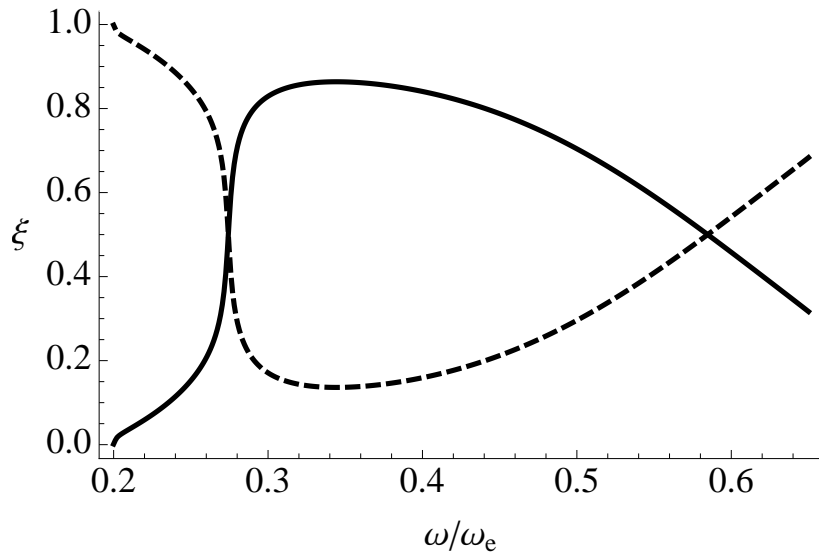


FIG. 6. Plots of the fraction of the total energy in the core (solid) and the cladding (dashed) as a function of frequency for the low-loss surface mode of the cylindrical metamaterial-dielectric waveguide.

The results above show that embedding three-level Λ atoms in the core of a metamaterial-dielectric waveguide provides coherently controllable slowing of signal field with reduced losses. In the next section, we discuss the loss mechanism in the metamaterial waveguide and explain how the existence of the low loss TM surface mode is due to Γ_m taking a value less than Γ_e .

IV. DISCUSSION

In our calculation, we have treated the claddings as if they are infinitely thick. In practice, the TM surface modes for the pump and signal fields only penetrate into the metamaterial to some finite depth, so a finite thickness of the cladding regions is sufficient. The skin depth for both the signal and pump fields is about $0.1\lambda_s$ with λ_s the wavelength of the signal at the EIT resonance. As typical fishnet metamaterial has thickness of $0.1\lambda_s$ [15], a few layers of fishnet metamaterial (with dielectric spacers in between) is sufficient for the cladding.

We find that the attenuation dip near ω_s (Fig. 5) is not as pronounced in the presence of transverse field confinement in the cylindrical waveguide. The flat metamaterial-dielectric interface has near-zero attenuation near this frequency, whereas the attenuation dip in the cylindrical guide does not approach zero. The cause for the non-zero attenuation is that the transverse confinement limits the amount of electromagnetic energy that is ejected from the lossy metamaterial.

To illustrate why transverse confinement lessens the attenuation dip near ω_s , consider a cylindrical guide in the limit of an infinite radius. In this limit, the interface that the surface wave propagates along becomes flat and the dispersion relation for the cylindrical waveguide reduces to that for a flat interface waveguide. In the case of the flat interface, the energy can reside almost completely on the dielectric side leading to near-zero loss [9]. Alternatively, for a cylindrical guide in the zero core radius limit, the core almost vanishes; hence, the energy propagates almost entirely in the lossy metamaterial, which leads to significant loss. For a waveguide with a finite radius, the fraction of the total energy that can be confined to the dielectric lies between 0 and 1; thus, there is a non-negligible amount of attenuation due to the field interacting with the metamaterial cladding.

We have considered fishnet metamaterial for our scheme but the scheme is not restricted to using fishnet metamaterial. Other types of metamaterial are applicable, as long as they have permittivity and permeability functions resembling equations (1) and (2), with $\epsilon'_{MM} < 0$, $0 < \mu'_{MM} < 1$ and $\Gamma_m < \Gamma_e$. Using other metamaterials is of interest because when layered fishnet metamaterials are bent to form a cylindrical cladding, the metamaterial may not interact with the fields isotropically. There now exist various promising fabrication techniques to building bulk metamaterials instead of layered ones [15, 23], which are isotropic.

We remark that our scheme is limited to narrowband signal pulses. This is because the spectral width of the EIT transparency window is related to the intensity of the pump field [10]. As slowing light requires a low intensity pump field, the resultant transparency window is spectrally narrow. The signal pulse must fit within the transparent spectral

window, so the signal pulse must be narrowband.

V. SUMMARY

We have considered two cylindrical dielectric-core waveguides of equal size and dimensions, with one having a metal cladding and the other a fishnet metamaterial cladding; other types of metamaterial cladding are also considerable. The permittivities of both claddings are given by the Drude model, but the metal cladding has constant permeability and the metamaterial cladding is given by the modified-Drude model. The signal field propagating through the waveguides are controllably slowed via electromagnetically induced transparency (EIT) due to the three-level Λ atoms homogeneously doped throughout the dielectric core of each waveguide. We have shown that the low-loss TM surface mode exists in the metamaterial waveguide provided that the permeability of the metamaterial cladding has a real part less than unity, i.e. $\text{Re}(\mu) < 1$, and the magnetic damping rate is less than the electric damping rate, i.e. $\Gamma_m < \Gamma_e$. Previous work has shown that this condition is always true for certain optical metamaterials [18]. We predict that such cylindrical metamaterial-cladded waveguide is capable of delivering the same slowing of light but with reduced loss compared to the cylindrical metal-cladded waveguide.

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